

# Massive higher spins and holography<sup>1</sup>

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**Abstract.** We review recent progress towards the understanding of higher spin gauge symmetry breaking in AdS space from a holographic vantage point. According to the AdS/CFT correspondence,  $\mathcal{N} = 4$  SYM theory at vanishing coupling constant should be dual to a theory in AdS which exhibits higher spin gauge symmetry enhancement. When the SYM coupling is non-zero, all but a handful of HS currents are violated by anomalies, and correspondingly local higher spin symmetry in the bulk gets spontaneously broken. In agreement with previous results and holographic expectations, we find that, barring one notable exception (spin 1 eating spin 0), the Goldstone modes responsible for HS symmetry breaking in AdS have non-vanishing mass even in the limit in which the gauge symmetry is restored. We show that spontaneous breaking à la Stückelberg implies that the mass of the relevant spin  $s' = s - 1$  Goldstone field is exactly the one predicted by the correspondence.

## 1. Introduction

The AdS/CFT correspondence [1] between IIB superstring theory on  $AdS_5 \times S^5$  with  $N$  units of 5-form flux and  $SU(N)$   $\mathcal{N} = 4$  SYM theory in  $d = 4$  is most often discussed in the limit of large AdS radius. In this limit the AdS side is under control, since the higher spin fields become extremely massive and decouple from the (super)gravity modes. Instead, in the CFT side this corresponds to the limit of large 't Hooft coupling. Therefore one can make predictions for the strongly coupled CFT using the correspondence but these can generally only be checked for certain protected objects.

More recently the opposite limit, in which the CFT is weakly coupled, has been discussed by a number of people [2]. In [3] the string spectrum in an  $AdS_5 \times S^5$  background at small radius was extrapolated and it precisely matches the operator spectrum of free  $\mathcal{N} = 4$  SYM in the planar limit. In particular the limit of zero YM coupling has been conjectured to be dual to a massless higher spin field theory which, although inconsistent when coupled to gravity in flat space-time, can be consistently defined in AdS spaces [4] (for a review see [5] and references therein). Turning on the coupling in the YM side corresponds in AdS to a Higgs mechanism, in which the massless higher spin fields develop a mass, essentially by eating lower spin Goldstone fields. This phenomenon was termed ‘La Grande Bouffe’ in [6]. The remaining massless fields will all be contained in the supergravity multiplet. In the dual CFT at zero coupling there

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are infinitely many higher spin conserved currents (in one-to-one correspondence with the AdS higher spin gauge fields). The CFT counter-part of ‘La Grande Bouffe’ is thus the anomalous violation of these conserved currents when the coupling is turned on, with the only remaining conserved currents lying in the energy momentum tensor multiplet.

The simplest example of a (bosonic) higher spin  $s$  field is that of a tensor with  $s$  completely symmetrised spacetime indices. For such an object, in flat spacetime, the massless limit of a massive spin  $s$  field gives rise to  $s + 1$  massless fields of spins  $0, 1, \dots s$ . However the AdS/CFT correspondence predicts that the massless limit of a massive spin  $s$  field in AdS is a massless spin  $s$  field and a massive spin  $s - 1$  field. The reason for this is that HS currents  $J_{i_1 \dots i_s}$  with  $s > 2$  occur in  $\mathcal{N} = 4$  SYM, where they are conserved at vanishing coupling  $g = 0$ , and conformal invariance fixes the dimension of such a spin  $s$  conserved current on the  $d$  dimensional boundary to be  $s + d - 2$ . Interactions are responsible for their anomalous violation

$$\partial^{i_1} J_{i_1 \dots i_s} = g \mathcal{X}_{i_2 \dots i_s} \quad ,$$

and in the zero coupling limit the dimension of  $\mathcal{X}$  is  $s + d - 1$ . This implies that  $\mathcal{X}$  is not a conserved spin  $s - 1$  current when  $g = 0$ , and therefore one expects it to be dual to a massive field in the bulk.

In a recent paper [7] the Stückelberg formulation of bosonic massive higher spin fields (with completely symmetrised spacetime indices) in AdS was derived (see also [8, 9] for similar results). These massive equations in AdS were then used to extrapolate the massless limit, and one indeed obtains a massless spin  $s$  field and a massive spin  $s - 1$  field in line with the CFT predictions (this phenomenon has also been discussed from a cosmological viewpoint in [8] where it was termed ‘partial masslessness’). The mass of the spin  $s - 1$  field one obtains in this way is precisely the one predicted by AdS/CFT.

Here we first review the results of [7]. We will only focus on higher spin bosons whose spacetime indices are completely symmetrised, although the fermionic case has been analysed in [10] and turns out to reveal similar features. We then conclude with some comments on how to generalize our analysis to other higher spin representations and some speculations on how to trigger *la Grande Bouffe* in the AdS bulk.

## 2. Higgs á la Stückelberg for higher spin fields: flat space vs AdS

An easy way to derive the Stückelberg formulation of a massive spin  $s$  field in flat  $D$  dimensional spacetime is to consider a massless spin  $s$  field in  $D + 1$  dimensions [11], described in terms of a symmetric rank  $s$  tensor<sup>2</sup>  $\Phi_{M_1 \dots M_s}$  ( $M_i$  are  $D + 1$ -dimensional spacetime indices) satisfying the condition

$$\Phi^L{}_{L'}{}^M{}_{M'} \Phi_{M_5 \dots M_s} = 0 \quad . \quad (1)$$

The resulting equation is invariant with respect to the gauge transformation

$$\delta \Phi_{M_1 \dots M_s} = s \partial_{(M_1} \epsilon_{M_2 \dots M_s)} \quad , \quad (2)$$

where the gauge parameter  $\epsilon$  is symmetric and traceless [12]. In order to perform a ‘massive’ KK dimensional reduction, one has to consider the various component fields to depend harmonically on the extra coordinate  $y$ ,

$$\Phi_{\mu_1 \dots \mu_{s-k} y \dots y}(x, y) = (i)^k \phi_{\mu_1 \dots \mu_{s-k}}^{(s-k)}(x) e^{imy} + \text{c.c.} \quad . \quad (3)$$

<sup>2</sup> Parentheses (...) denote symmetrization of spacetime indices with strength one. This is the origin of factors of  $s$  in several formulae.

By taking linear combinations one can choose the fields  $\phi^{(s-k)}$  to be real in  $D$  dimensions. The gauge transformation (2) becomes for the  $D$ -dimensional fields

$$\delta\phi_{\mu_1\ldots\mu_{s-k}}^{(s-k)} = (s-k)\partial_{(\mu_1}\epsilon_{\mu_2\ldots\mu_{s-k})}^{(s-k-1)} + km\epsilon_{\mu_1\ldots\mu_{s-k}}^{(s-k)} \quad , \quad (4)$$

where the  $D$ -dimensional gauge parameters related to the  $(D+1)$ -dimensional ones by means of

$$\epsilon_{\mu_1\ldots\mu_{s-k-1}y\ldots y}(x, y) = (i)^k \epsilon_{\mu_1\ldots\mu_{s-k-1}}^{(s-k-1)} e^{imy} + \text{c.c.} \quad . \quad (5)$$

If  $m \neq 0$ , from eq. (4) it turns that only  $\phi^{(s)}$  does not transform algebraically with respect to any gauge transformations. Actually, not all the lower spin components can be put to zero fixing their gauge invariance, because of the traceless constraint on the gauge parameters. The remaining lower spin fields, that are identically zero on shell, are the auxiliary fields of the massive theory [13], and one ends up with an equation for a massive spin  $s$  field  $\phi^{(s)}$ . If  $m = 0$ , instead, none of the gauge parameters can be used to gauge away any of the fields, and therefore all the fields  $\phi^{(s-k)}$ ,  $k = 0, 1, \dots, s$ , become massless.

We now want to consider the same system of equations in AdS. In particular, we consider the field equation for  $\phi^{(s-1)}$ , and we gauge away  $\phi^{(s-2)}$  and  $\phi^{(s-3)}$  using  $\epsilon^{(s-2)}$  and  $\epsilon^{(s-3)}$ . This implies that only the fields  $\phi^{(s)}$  and  $\phi^{(s-1)}$  will appear in the equation, while all the other fields  $\phi^{(s-k)}$ , with  $k = 4, \dots, s$  are auxiliary fields. The only gauge invariance left is the one with respect to the *traceless* gauge parameter  $\epsilon^{(s-1)}$ , and in AdS it will require the addition of a mass term for  $\phi^{(s-1)}$ . We then go to the  $m = 0$  limit, and we check whether the mass term we included leads to a gauge symmetry enhancement or not. Since we don't find any inconsistency, that would indicate gauge symmetry enhancement, we can continue the procedure to  $m = 0$  and we end up with a massless spin  $s$  and a massive spin  $s - 1$ .

Before we proceed, we first observe that gauge invariance (*i.e.* masslessness) in AdS implies the presence of a mass-like term in the field equations, proportional to the inverse of the AdS radius  $L$ . In particular, the equation for a spin  $l$  field

$$\begin{aligned} & \square\phi_{\mu_1\ldots\mu_l}^{(l)} - l\nabla_{(\mu_1}(\nabla \cdot \phi^{(l)})_{\mu_2\ldots\mu_l)} + \frac{l(l-1)}{2}\nabla_{(\mu_1}\nabla_{\mu_2}\phi^{(l)\lambda}_{\lambda\mu_3\ldots\mu_l)} \\ & - M_{AdS}^2 \phi_{\mu_1\ldots\mu_l}^{(l)} - \tilde{M}_{AdS}^2 g_{(\mu_1\mu_2}\phi^{(l)\lambda}_{\lambda\mu_3\ldots\mu_l)} = 0 \quad , \end{aligned} \quad (6)$$

with

$$M_{AdS}^2 = \frac{(l-2)(D-1) + (l-1)(l-4)}{L^2} \quad , \quad \tilde{M}_{AdS}^2 = \frac{l(l-1)}{L^2} \quad , \quad (7)$$

is gauge invariant with respect to

$$\delta\phi_{\mu_1\ldots\mu_l}^{(l)} = l\nabla_{(\mu_1}\epsilon_{\mu_2\ldots\mu_l)}^{(l-1)} \quad , \quad (8)$$

where  $\epsilon$  is traceless.

We now consider in AdS the equation for the spin  $s - 1$  Stückelberg field  $\phi^{(s-1)}$ , that we denote here by  $\chi^{(s-1)}$  for clarity. Requiring that this equation is invariant with respect to the gauge transformations

$$\delta\phi_{\mu_1\ldots\mu_s}^{(s)} = s\nabla_{(\mu_1}\epsilon_{\mu_2\ldots\mu_s)} \quad , \quad \delta\chi_{\mu_1\ldots\mu_{s-1}}^{(s-1)} = m\epsilon_{\mu_1\ldots\mu_{s-1}} \quad , \quad (9)$$

with  $\epsilon$  traceless, implies the presence of a mass proportional to the inverse of the AdS radius, so that one ends up with the equation

$$\square\chi_{\mu_1\ldots\mu_{s-1}}^{(s-1)} - (s-1)\nabla_{(\mu_1}(\nabla \cdot \chi^{(s-1)})_{\mu_2\ldots\mu_{s-1})}$$

$$\begin{aligned}
& + \frac{(s-1)(s-2)}{2} \nabla_{(\mu_1} \nabla_{\mu_2} \chi^{(s-1)\lambda}_{\lambda\mu_3\ldots\mu_{s-1})} \\
& - m(\nabla \cdot \phi^{(s)})_{\mu_1\ldots\mu_{s-1}} + (s-1)m \nabla_{(\mu_1} \phi^{(s)\lambda}_{\lambda\mu_2\ldots\mu_{s-1})} \\
& - \frac{(s-1)[(D-1) + (s-2)]}{L^2} \chi^{(s-1)}_{\mu_1\ldots\mu_{s-1}} = 0 \quad .
\end{aligned} \tag{10}$$

We thus would like to compare this mass term with the first of eqs. (7), where  $l = s-1$ . They are definitely different, which means that no symmetry enhancement occurs when  $m = 0$ , and any massive spin  $s$  field in the limit of zero mass decomposes into a massless spin  $s$  field and a massive spin  $s-1$  field. In other words, the new feature of AdS is the fact that the auxiliary field structure is preserved for the spin  $s-1$  field even when the spin  $s$  field becomes massless. Note that our procedure leaves undetermined a possible mass term of the form  $1/L^2 g_{(\mu_1\mu_2} \chi^{(s-1)\lambda}_{\lambda\mu_3\ldots\mu_s)}$  in eq. (10), since  $\chi^{(s-1)\lambda}_{\lambda}$  is gauge invariant with respect to (9). This is not an issue as long as we focus on the field equations, since  $\chi^{(s-1)\lambda}_{\lambda}$  can be put to zero on shell using the lower rank equations. Nevertheless, the whole set of equations can be derived from a lagrangian once the correct equations for the auxiliary fields are introduced, in a similar way to the flat space case.

The difference between the AdS mass term and this mass term (for simplicity we define  $s' = s-1$  from now on) is

$$- \frac{2(D-1) + 4(s'-1)}{L^2} \chi^{(s')}_{(\mu_1\ldots\mu'_{s'})} \quad . \tag{11}$$

We therefore get

$$M^2 L^2 = 2(D-1) + 4(s'-1) \quad . \tag{12}$$

In  $D = 5$  ( $d = 4$ ) this equation becomes

$$M^2 L^2 = 4(s' + 1) \quad . \tag{13}$$

This is exactly what we get from the standard relation between mass in AdS and dimension of the dual operator in the boundary theory [14, 6],

$$M^2 L^2 = \Delta(\Delta - 4) - \Delta_{min}(\Delta_{min} - 4) \quad , \tag{14}$$

with

$$\Delta = s' + 4 \quad , \quad \Delta_{min} = s' + 2 \quad , \tag{15}$$

which is exactly the dimension of the corresponding spin  $s'$  operator at vanishing Yang-Mills coupling. For arbitrary dimension  $d = D-1$ ,  $\Delta_{min} = s' + d - 2$  represents the unitary bound for the dimension of a spin  $s'$  current, *i.e.* a totally symmetric rank  $s'$  classically conserved tensor current, and the identity (14), with 4 substituted with  $d$ , is satisfied with  $\Delta = s' + d$ .

The case of an anomalous spin  $s = 1$  (axial) vector current, for which the relevant Stückelberg field in the bulk is a massless (pseudo) scalar, is special. Indeed, eqs. (6,7) are meaningless for a scalar field, while eq. (10) shows that for  $s = 1$  the mass for  $\chi^{(0)}$  vanishes. This agrees with the mass/dimension relation for scalar fields,

$$M^2 L^2 = \Delta(\Delta - d) \quad , \tag{16}$$

since  $\chi^{(0)}$  is dual to a naively marginal scalar operator of dimension  $\Delta = d = D-1$ .

### 3. Conclusions and perspectives

We would like to conclude with some comments on how to generalize our analysis to other higher spin representations and some speculations on how to trigger *la Grande Bouffe* in the AdS bulk.

In the introduction we already mentioned how the case of a spontaneously broken fermionic spin  $s + 1/2$  symmetry in AdS can be described [10] along the lines of the bosonic spin  $s$  totally symmetric tensors [7] reviewed above. The two cases can indeed be related by exploiting the global supersymmetries present in AdS even after HS symmetry breaking. By the same token one should be able to relate more general HS bosonic and fermionic representations that appear in the AdS HS supermultiplets for  $D = 5$  and higher. Indeed the analysis of [22, 23, 24] shows that all the relevant Goldstone modes are present in the free SYM spectrum that are needed to achieve HS breaking in a way compatible with full  $PSU(2, 2|4)$  symmetry. In particular, it is shown there that the string spectrum on  $AdS_5 \times S^5$  can be extrapolated to the HS enhancement point where it can be precisely matched with the spectrum of single trace gauge invariant operators and decomposed into representations of  $HS(2, 2|4)$ . The latter, indicated as *YT-pletons*, are in one to one correspondence with the Yang Tableaux (YT) compatible with the cyclicity of the trace over the colour indices which is the counterpart of ‘level matching’, *i.e.* the closure of the string. To wit, the HS gauge fields belong to the *doubleton*  $\Phi_2$  corresponding to the bilinear gauge invariant operators  $\mathcal{O}_2^{AB} = \text{tr}(W^A W^B)$ , where  $W^A$  denotes any of the ‘letters’ of the SYM ‘alphabet’, *i.e.* any of the fundamental SYM fields or derivatives thereof, modulo field equations. Similarly there are two *tripletons*: the *f-triplet*  $\mathcal{X}_3$ , corresponding to the totally antisymmetric trilinear operators  $\mathcal{O}_{3,f}^{[ABC]} = \text{tr}(W^A [W^B, W^C])$ , expressible in terms of the structure constants  $f_{abc}$ , which contains the first generation of Goldstone modes, and the *d-triplet*  $\Phi_3$ , corresponding to the totally symmetric trilinear operators  $\mathcal{O}_{3,d}^{(ABC)} = \text{tr}(W^A \{W^B, W^C\})$ , expressible in terms of the cubic Casimir  $d_{abc}$ , which accounts for the first KK recurrence of the supergravity fields and their HS cousins. No *hooked* YT with three boxes are compatible with gauge symmetry, *i.e.* cyclicity of the trace or equivalently level matching / closure of the string. By the same token no two-particle *tripletons* are allowed that would correspond to double-trace trilinear operators. Double-trace operators first appear in *tetrapletons*, *e.g.*  $\Phi_{2,2} = \Phi_2 \Phi_2$ . Other *tetrapletons* can be of various kinds. The *q-tetraplet*  $\Phi_4$ , corresponding to the totally symmetric quadrilinear operators  $\mathcal{O}_{4,q}^{(ABCD)} = \text{tr}(\{W^A, W^{(B}\{W^C, W^D\}\})$ , expressible in terms of the quartic Casimir  $q_{abcd} = \delta^{ef} d_{e(ab} d_{cd)f}$ , which accounts for the second KK recurrence of the supergravity fields and their HS cousins. The *w-tetraplet* (*w* stands for window:  $w \approx ff$ )  $\mathcal{X}_4$ , corresponding to the quadrilinear operators  $\mathcal{O}_{4,w}^{[AB][CD]} = \text{tr}([W^A, W^B][W^C, W^D])$ , expressible in terms of  $w_{[ab][cd]} = \delta^{ef} f_{abe} f_{cdf}$ , which accounts for the second generation of Goldstone modes. Moreover, the *h-tetraplet* (*h* stands for hooked)  $\Psi_4$ , corresponding to the mixed symmetry quadrilinear operators  $\mathcal{O}_{4,h}^{(A[B]C]D]} = \text{tr}(\{W^A, W^{[B}\{W^C, W^D]\})$ , expressible in terms of  $h_{(a[b]cd]} = \delta^{ef} d_{ea[b} f_{cd]f}$ , which accounts for genuinely massive HS fields.

We are now ready to speculate on the structure of the Lagrangian, if any, that governs the dynamics of the massless and massive HS multiplets at the point of HS enhancement. We heavily rely on the known properties of free SYM theory that should provide the holographic dual. We start by choosing a normalization such that two-point functions are normalized to 1 and focus only on the master fields<sup>3</sup>  $\Phi_2$ ,  $\mathcal{X}_3$ ,  $\mathcal{X}_4$ , then schematically

$$\begin{aligned} \mathcal{L}_{HS} = & \Phi_2 \Phi_2 + \frac{1}{N} \Phi_2 \Phi_2 \Phi_2 + \frac{1}{N^2} \Phi_2 \Phi_2 \Phi_2 \Phi_2 + \dots \\ & \mathcal{X}_3 \mathcal{X}_3 + \frac{1}{N} \mathcal{X}_3 \mathcal{X}_3 \Phi_2 + \frac{1}{N^2} (\mathcal{X}_3 \mathcal{X}_3 \Phi_2 \Phi_2 + \mathcal{X}_3 \mathcal{X}_3 \mathcal{X}_3 \mathcal{X}_3) + \dots \end{aligned} \quad (17)$$

<sup>3</sup> We are using the term master field in a loose sense. Strictly speaking the massless *doubleton*  $\Phi_2$  requires two master fields in Vasiliev description [4, 5], a master connection  $\mathcal{W}$  and a master scalar curvature  $\mathcal{R}$ . For massive HS multiplets such as  $\mathcal{X}_3$  and  $\mathcal{X}_4$  no explicit master field description is available at present.

$$\mathcal{X}_4\mathcal{X}_4 + \frac{1}{N}(\mathcal{X}_4\mathcal{X}_3\mathcal{X}_3 + \mathcal{X}_4\Phi_2\Phi_2) + \frac{1}{N^2}(\mathcal{X}_4\mathcal{X}_3\mathcal{X}_3\Phi_2 + \mathcal{X}_4\Phi_2\Phi_2\Phi_2 + \mathcal{X}_4\mathcal{X}_4\Phi_2\Phi_2 + \mathcal{X}_4\mathcal{X}_4\mathcal{X}_3\mathcal{X}_3) + \mathcal{X}_4\mathcal{X}_4\mathcal{X}_4\mathcal{X}_4) + \dots$$

Notice that any term in  $\mathcal{L}_{HS}$  encompasses an infinite number of terms / couplings involving the component HS fields as well as their derivatives to arbitrary high order. The schematic form of  $\mathcal{L}_{HS}$  seems already to be in conflict with the consistent truncation to the massless HS theory that only involves  $\Phi_2$  [15], and sets  $\mathcal{X}_3, \mathcal{X}_4$  to zero. Notice however that the offending coupling  $\mathcal{X}_4\Phi_2\Phi_2$  is extremal and as such, in the spirit of the holographic correspondence, is expected to arise from a purely boundary contribution since the relevant bulk integral would diverge (at least for the lowest scalar components and the rest should follow from  $HS(2, 2|4)$  if not  $PSU(2, 2|4)$  symmetry). Alternatively one can define  $\mathcal{X}_4$  in such a way that it is ‘orthogonal’ to  $\Phi_{2,2} = \Phi_2\Phi_2$ . From now on we assume that this is the case. It is also convenient to rescale the master fields  $\Phi_2, \mathcal{X}_3, \mathcal{X}_4$  by a factor of  $N$  so that  $\mathcal{L}_{HS}$  displays an overall factor of  $N^2$  (that should in the end become an  $N^2 - 1$  for  $SU(N)$ )

$$\begin{aligned} \mathcal{L}_{HS} = & N^2(\Phi_2\Phi_2 + \Phi_2\Phi_2\Phi_2 + \Phi_2\Phi_2\Phi_2\Phi_2 + \dots \\ & \mathcal{X}_3\mathcal{X}_3 + \mathcal{X}_3\mathcal{X}_3\Phi_2 + \mathcal{X}_3\mathcal{X}_3\Phi_2\Phi_2 + \mathcal{X}_3\mathcal{X}_3\mathcal{X}_3\mathcal{X}_3(1 + \mathcal{O}(1/N^2)) + \dots \\ & \mathcal{X}_4\mathcal{X}_4 + \mathcal{X}_4\mathcal{X}_3\mathcal{X}_3 + \mathcal{X}_4\Phi_2\Phi_2 + \mathcal{X}_4\mathcal{X}_3\mathcal{X}_3\Phi_2 + \mathcal{X}_4\Phi_2\Phi_2\Phi_2 + \mathcal{X}_4\mathcal{X}_4\Phi_2\Phi_2 + \\ & + \mathcal{X}_4\mathcal{X}_4\mathcal{X}_3\mathcal{X}_3(1 + \mathcal{O}(1/N^2)) + \mathcal{X}_4\mathcal{X}_4\mathcal{X}_4\mathcal{X}_4(1 + \mathcal{O}(1/N^2)) + \dots \end{aligned} \quad (18)$$

As indicated non-planar corrections (subleading in  $1/N$ ) affect some of the couplings, *i.e.* those involving at least six  $f$ ’s. Inclusion of the other *YT-pletons* should not change the structure in any significant way.

Assuming that HS symmetry be strong enough to fix completely the detailed structure of  $\mathcal{L}_{HS}(\Phi_2, \mathcal{X}_3, \mathcal{X}_4, \dots)$ , we can take it as the starting point to discuss HS symmetry breaking. Holography suggests there must be a massless (complex) scalar, the (complexified) dilaton, dual to the SYM interaction lagrangian  $\mathcal{L}_{SYM}$ , that could acquire a VEV preserving exact  $PSU(2, 2|4)$ , while breaking  $HS(2, 2|4)$ . This scalar<sup>4</sup> is a combination of the the massless scalar singlets that appear in  $\Phi_2, \mathcal{X}_3$  and  $\mathcal{X}_4$ . Once it takes a VEV, all but a handful of HS gauge fields in  $\Phi_2$  become massive by eating lower spin Goldstone modes in  $\mathcal{X}_3$  which in turn absorb the second generation Goldstone modes in  $\mathcal{X}_4$ . In order to make contact with the discussion of the totally symmetric bosonic spin  $s$  tensors, the parameter  $M$  there should be thought of as the relevant VEV.

Admittedly the resulting picture is still obscure due to the lack of knowledge of an efficient description of massive HS multiplets á la Vasiliev. Yet it is quite appealing how subtle effects such as the anomalous violation of HS currents might be described in rather compact terms as hinted at above. To be concrete, one may try to use HS symmetry, that acts linearly on the master fields prior to its spontaneous breaking, to fix the relative strength of some of the trilinear couplings that are responsible for the spontaneous breaking and the consequent mass generation once the ‘dilaton’ gets a VEV. This is our main goal in the near future.

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<sup>4</sup> We here mean the component dual to the SYM coupling  $g_{YM}$ , which is known to be exactly marginal. We neglect for the time being  $\vartheta_{YM}$  that is also exactly marginal but can only show up in non-perturbative corrections induced by (D)-instantons[25] .

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